Screening Effects in Superfluid Nuclear and Neutron Matter within Brueckner Theory

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Effects of medium polarization are studied for $^1S_0$ pairing in neutron and nuclear matter. The screening potential is calculated in the RPA limit, suitably renormalized to cure the low density mechanical instability of nuclear matter. The selfenergy corrections are consistently included resulting in a strong depletion of the Fermi surface. All medium effects are calculated based on the Brueckner theory. The $^1S_0$ gap is determined from the generalized gap equation. The selfenergy corrections always lead to a quenching of the gap, which is enhanced by the screening effect of the pairing potential in neutron matter, whereas it is almost completely compensated by the antiscreening effect in nuclear matter.

I. INTRODUCTION

A satisfactory description of superfluidity in nuclear matter has not yet been achieved despite almost fifty years of research have elapsed since the first application of the BCS theory to nuclear systems [1]. Somewhat at variance with the electron pairing in superconductors the pairing in nuclear systems results from the interplay between the direct action of the bare nuclear force and the action induced by the medium polarization. The attractive components of the bare nuclear interaction have led to the investigation of several pairing configurations, e.g. neutron-neutron or proton-proton pairing in the $^1S_0$ channel in neutron stars [2] disregarding possible repulsive effect exerted by screening of the force via the medium. A pairing suppression has in fact been found by most calculations of pairing in neutron matter (see Ref. [3] and references therein). On the contrary, other pairing configurations have not yet been explored since the repulsive components of the direct nuclear interaction cannot support the formation of Cooper pairs. But there are strong indications that, in a nuclear rather than neutron matter environment, the medium polarization of the interaction can favor the formation of Cooper pairs similar to the lattice vibrations in ordinary superconductors. These indications come both from nuclear matter calculations and from finite nuclei. In nuclear matter the medium enhancement of neutron-neutron $^1S_0$ pairing is to be traced back to the proton particle-hole excitations [4], and in finite nuclei to the surface vibrations [5].

Another distinctive feature of the nuclear environment is the presence of strong short range correlations that induce two effects relevant for the pairing: one is the depletion of the Fermi surface, which is experimentally supported by measurements of electron scattering on ^{208}Pb [6], the other one is the strong mass renormalization caused by short-range particle-particle correlations [7].

The two effect conspire against the pairing formation:

The depletion of the Fermi sea reduces the phase space available for particle-particle virtual transitions around the Fermi surface, the mass renormalization enhances the dispersive effect of the mean field [8].

Therefore a complete microscopic treatment of the very subtle pairing problem requires vertex and selfenergy corrections to be treated and to be considered on the same footing. In a previous paper[4] we made a study of these in-medium effects under several simplifying assumptions. First came the approximation to replace the Born term of the pair interaction in the S=0, T=1 channel by the Gogny force [9]. Though in that channel the Gogny force is not dissimilar to the action of the bare force (see Ref. [10]), it shows a little too much attraction for momenta characterizing saturation. The first improvement in the present work is then the use of a realistic two body force (V18 [11], see below) in the Born term. Secondly we will use as vertices in the induced interaction a force which is based on a more modern G-matrix calculation [12] as this was the case for the Gogny force [9]. Thirdly we corrected an unfortunate phase error which slipped into the evaluation of the induced force at least for the symmetric nuclear matter case in the S=0, T=1 channel what gave raise to a too strong anti-screening effect. With these corrections and improvements we now get reasonable renormalization effects of the pairing force and we calculate the corresponding gaps as a function of density in pure neutron matter as well as in symmetric nuclear matter.

In detail the paper is organized as follows. In Sec. II the generalized gap equation is reviewed along with the approximations on the pairing potential and selfenergy, which lead to the determination of the energy gap. In Sec. III the screening interaction is discussed in the RPA limit, and then the summation of bubble diagrams and the resummation of dressed bubble diagrams both using the Landau parameters are derived. In Sec. IV the results are presented: first, for the separate contributions of particle-hole (ph) scalar, vector and isovector excitations in neutron matter and nuclear

matter; second, the solution of the gap equation for $^1S_0$ pairing with a discussion of the effects of the selfenergy corrections and medium polarization potential. Section V is devoted to the comparison with other calculations and to the conclusions.

II. GENERALIZED GAP EQUATION

The spectrum of a superfluid homogeneous Fermi system is derived from the generalized gap equation[13, 14, 15]:

$$\Delta_{\boldsymbol{k}}(\omega) = \sum_{k'} \int \frac{d\omega'}{2\pi i} \mathcal{V}_{\boldsymbol{k},\boldsymbol{k}'}(\omega,\omega') F_{\boldsymbol{k}'}(\omega'), \qquad (1)$$

where $\mathcal{V}$ is the sum of all irreducible NN interaction diagrams and $F_k(\omega)$ is the anomalous propagator. The class of diagrams selected for the present calculation is plotted in Fig. 1.

In nuclear matter $\mathcal{V}$ can be approximated by the bare interaction $\mathcal{V}_0$(diagram (a)), which is responsible of the formation of Cooper pairs, and the class of bubble insertions, which play the role of screening. In turn, the screening interaction $\mathcal{V}_1$ can be split into two parts: the one-bubble term (diagram (b)) containing only the mixed configuration with particle-particle (pp) plus ph excitations, the multi-bubble term also containing all insertions of pure ph interaction vertices (diagram (c)). The splitting is a convenient way to point out that mixed vertices and pure ph vertices have to be treated on different footing, as discussed in Ref. [4]. The first bubble diagram can also be seen as the lowest order correction to the Born term. As discussed afterwards, the vertex insertions in diagrams (b) and (c) are described by a Brueckner G-matrix.
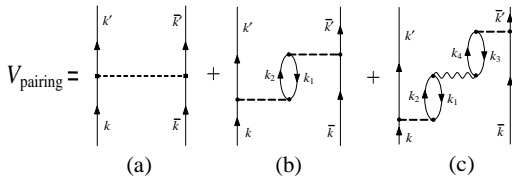

FIG. 1: Pairing interaction with screening in the RPA approximation. The short-dashed line represents the bare interaction, long-dashed lines the G-matrix, the wiggly line the p-h residual interaction resummed to all orders. All vertices are to be understood as anti-symmetrised matrix elements. The notations are: $k = (\vec{k}, \sigma, \tau), \bar{k} = (-\vec{k}, -\sigma, \tau)$.

Within the Brueckner theory of nuclear matter, the single particle (sp) energy spectrum of the non superfluid state is derived from the hole-line expansion of the mass operator. The BHF approximation, extended to include depletion of the Fermi surface due to the strong ground state correlations, gives an important quenching of the pairing gap and hence it can not be neglected. The pp correlations on the pairing interaction are embodied in the gap equation itself [1]. Consistently, the interaction vertices in the screening term should be described in terms of the G-matrix. In addition, the lat-

ter must be dressed according to the Babu-Brown theory of the induced interaction [16, 17] to avoid the low density instability problem of nuclear matter, as discussed later. Since the exact resummation of the bubble series (Bethe-Salpeter equation) with G-matrix is a prohibitive task, the ph vertex insertions can be conveniently evaluated in the Landau limit, and eventually replaced by the Landau parameters.

The effects of the selfenergy corrections have intensely been studied; in particular the depletion of the Fermi surface is expected to hinder the virtual transitions around the Fermi surface and thus its effect is to weaken the pairing correlations. It is of particular interest to understand to what extent the resulting quenching of pairing gap is compensated by in medium vertex corrections that in nuclear matter are strongly attractive. On the other hand, in the case of neutron matter selfenergy effects enhance quenching due to screening at variance with the predictions of recent Monte Carlo many body calculations [18].

Going beyond the pure BHF approximation, the main dispersive corrections arise from energy dependence of the self-energy, as shown in previous papers. The correction to the non superfluid propagator $G_k(\omega)$ is simply a renormalization factor of its pole part. This factor, named Z-factor, is

$$\mathcal{Z}^{-1} = 1 - [\frac{\partial \Sigma_k(\omega)}{\partial \omega}]_{\omega=\omega_k}, \qquad (2)$$

which measures the discontinuity of the occupation probability around the Fermi energy. Correspondingly, the abnormal propagator appearing in the gap equation (Eq. (1))

$$F_{\boldsymbol{k}}(\omega) = \frac{\Delta_k(\omega)}{G_k^{-1}(\omega)G_k^{-1}(-\omega) + \Delta_k^2(\omega)}, \qquad (3)$$

is renormalized by a factor Z^2. For a static interaction the gap function is also independent of energy and the energy integration can be performed analytically. Since the analytical structure of the abnormal propagator is not modified, the gap equation takes the same form as in the pure BCS case. One easily obtains:

$$\Delta_k = -\frac{1}{2} \int d^3\vec{k}' \mathcal{V}_{k,k'} \frac{Z_k Z_{k'} \Delta_{k'}}{\sqrt{(\varepsilon_{k'} - \varepsilon_F)^2 + \Delta_{k'}^2}}, \qquad (4)$$

where ε_F is the Fermi energy and ε_k is the on-shell self-energy. The preceding gap equation is equivalent to the BCS version except for the Z^2 factor which is modeling the effect of the interaction around the Fermi surface. Since the value of the Z-factor, though depending on the ground state correlations, is always less than unity in the vicinity of the Fermi surface, inevitably the energy gap will turn out quenched in this respect. Our predictions of the gap in nuclear and neutron matter presented in this paper rely on the solution of the latter equation.

III. SCREENING INTERACTION

A. One-bubble screening interaction

In a previous work the calculation of the screening interaction was simplified by using the Gogny force, which in fact reproduces most of the properties of a G-matrix.

In the present calculation we adopt the G-matrix itself and we try to reduce its complexity with reasonable approximations. For the sake of application to the pairing in the $^1S_0$ channel we select the two particle state with total spin S=0 and isospin T=1. Then the one-bubble interaction can be written as

$$< 1\bar{1}|\mathcal{V}_1|1'\bar{1}' > = \frac{1}{4}\sum_{2,2'}\sum_{ST}(-)^S(2S+1) < 12|G^{ph}_{ST}|1'2' >_A < 2'\bar{1}|G^{ph}_{ST}|2\bar{1}' >_A \Lambda^0(22'), \tag{5}$$

where $1 \equiv (\vec{k}_1, \sigma_1, \tau_1)$ $(1' \equiv (\vec{k}_{1'}, \sigma_{1'}, \tau_{1'}))$ and $\bar{1} \equiv (-\vec{k}_1, \sigma_1, \tau_1)$ $(\bar{1}' \equiv (-\vec{k}_{1'}, \sigma_{1'}, \tau_{1'}))$ are the momenta of the pair in the entrance (exit) channel. Λ is the static polarization part. The G-matrix is converted into the

ph sector, as it is required to solve the Bethe-Salpeter equation and to sum up the bubble series $\tilde{V}_2$. The standard recoupling procedure from pp sector to ph sector yields

$$G^{ph}_{ST} = \sum_c (2S_c+1)(2T_c+1)(-1)^{S_c+T_c} \left\{ \begin{array}{ccc} \frac{1}{2} & \frac{1}{2} & S_c \\ \frac{1}{2} & \frac{1}{2} & S \end{array} \right\} \left\{ \begin{array}{ccc} \frac{1}{2} & \frac{1}{2} & T_c \\ \frac{1}{2} & \frac{1}{2} & T \end{array} \right\} G_{S_c T_c}, \tag{6}$$

where the brackets are the 6j symbols. The sum runs over the spin S_c and isospin T_c of the pp channels included in the calculation. Since the G-matrix incorporates short range pp correlations, its momentum range is shrunk remarkably in comparison with the bare interaction, as shown in Fig. 2. At variance with the bare interaction, the the G-matrix cannot sustain large momentum transfers $\vec{q} = \vec{k} - \vec{k}'$, that justifies the approximation to average it around the Fermi surface, in the limit q=0. As a consequence the q dependence is only located in the integral of the polarization part, giving the Lindhard function

$$\sum_{\vec{k}} \Lambda^0_{\vec{k}, \vec{k}-\vec{q}} = \frac{N(0)}{g} \frac{1}{2} \left[-1 + \frac{1}{q}(1 - \frac{q^2}{4})ln \left| \frac{1-q/2}{1+q/2} \right| \right], \tag{7}$$

where g is the degeneracy parameter.

The two external vertices mixing pp and ph lines shown in Fig. 1 (b) in principle induce unlimited excitations in momentum space. But, since high momenta transitions are incorporated into G-matrix used as vertex interaction, the main contribution of the two-bubble diagram is concentrated in a domain as short as $2\ fm^{-1}$. Elsewhere the remaining contribution can neglected with respect to the bare interaction as shown in Fig. 2.

The problems with the calculation of the ph multi bubble contribution (diagram (c) in Fig. 1) are the following. First, the bubble series with G-matrix insertions have to be previously summed up. But, since
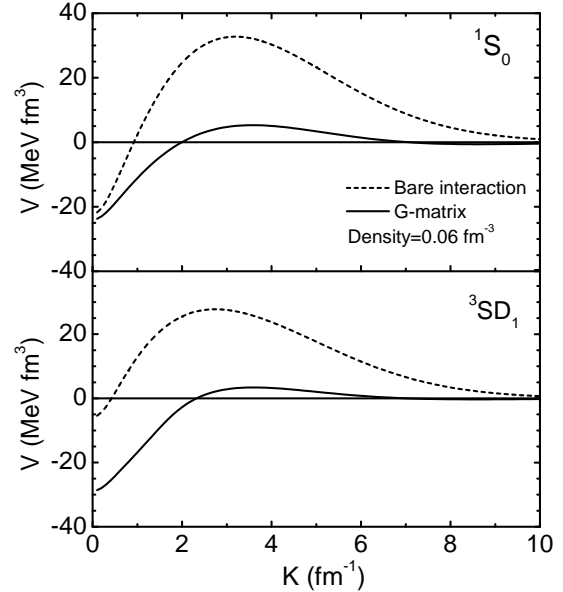

FIG. 2: Shrinking of G-matrix in momentum space.

the interaction vertices in the ph channel involve particle excitations around the Fermi surface, they can be approximated by the Landau parameters. Second, even replacing the bare interaction vertices by G-matrices, there appears the long low density singularity in the RPA in nuclear matter ($F_0 = -1$). This problem, discussed in Ref. [3] (see also references therein) is remedied by dressing the vertex insertions according to the

Babu-Brown induced interaction theory[16].

B. Landau parameters from the BHF approximation

The microscopic basis of the ph effective interaction can be set in terms of the energy functional of symmetric nuclear matter

$$N(0)f_{\sigma\tau,\sigma'\tau'}(\vec{k},\vec{k}') = \frac{\delta^2 E}{\delta n_{\sigma\tau}(\vec{k})\delta n_{\sigma'\tau'}(\vec{k}')} \quad (8)$$
$$= F + F'(\tau\cdot\tau') + G(\sigma\cdot\sigma') + G'(\sigma\cdot\sigma')(\tau\cdot\tau'),$$

where the density of states $N(0)$ is introduced to make the Landau parameters F, F', G and G' dimensionless. In BHF [7] approximation the energy functional is given by

$$E = \sum_k \frac{\hbar^2 k^2}{2m} + \frac{1}{2}\sum_{k_1,k_2}\langle k_1, k_2|G(\omega)|k_1, k_2\rangle_A, \quad (9)$$

where the subscript A means that the matrix element of the G-matrix is antisymmetrized. The index k stands for $\vec{k}$, σ and τ, momentum, spin and isospin, respectively. The G-matrix is understood to be calculated on the energy shell: $\omega = \epsilon_{k_1} + \epsilon_{k_2}$. The single particle energies are determined iteratively along with the G-matrix within the Brueckner selfconsistent scheme. One can determine the Landau parameters from the microscopic Brueckner theory in the BHF approximation, performing the double variational derivative, Eq. (8), of the energy per particle, Eq. (9). So doing, a number of contributions are generated that can be calculated one by one [19] in some approximation due to the complex structure of G-matrix. A simple and powerful way to calculate the Landau parameters is to suitably fit the BHF energy and the corresponding sp spectrum with a functional of the occupation numbers and then to perform the double derivative. A Skyrme-like functional has proved to reproduce accurately the equation of state (EoS) of symmetric as well as spin and isospin asymmetric nuclear matter [20]. Therefore we determine the Landau parameters in that way. A limitation of this procedure is that only a few partial wave components can be calculated, but for the purpose of the present investigation we only need the zero order Landau parameters. The latter are plotted in Fig. 3 as a function of the Fermi momentum. As expected F_0 exhibits the well known instability below the saturation point, which makes the RPA series difficult to handle. As in previous papers [4, 21] this drawback can be overcome by the induced interaction theory of Babu and Brown [16]. Leaving aside a description of this theory (see Refs. [17, 22]), we schematically write down the equation defining the ph induced interaction as follows

$$V_{ph} = V_d + V_{RPA}(V_{ph}). \quad (10)$$

The first term (direct term) is the BHF ph residual interaction, which, in the first order, is represented by the G-matrix. The second one (induced term) is the RPA bubble summation, in which the vertex insertions are given by V_{ph} itself instead of the direct term. The solution of the latter equation is quite simple if we replace the true V_d projection in the ph channel (ST) with the corresponding Landau parameter, since the way we extract the Landau parameters the direct term contains not only the effect of the G-matrix but also the rearrangement diagrams [22]. The numerical results are depicted in Fig. 3. The salient feature of the induced interaction is that the renormalization of F_0 prevents any singular behavior to occur below the saturation density. Otherwise the values do not differ from the Gogny interaction except for the high density behavior where the effect of the three body force makes F_0 much more repulsive than the Gogny force [4].

Therefore we dressed the residual interaction first with the short range correlations (G-matrix instead of bare interaction) and then by the renormalized long range correlations V_{ph} replacing the G-matrix in the RPA series. Since the calculation of the induced interaction with the G-matrix is a quite complex job, we have simplified the problem replacing the G-matrix with the Landau parameters. The way we determine the Landau parameters, the approximation turns out to be better than starting from the G-matrix itself.
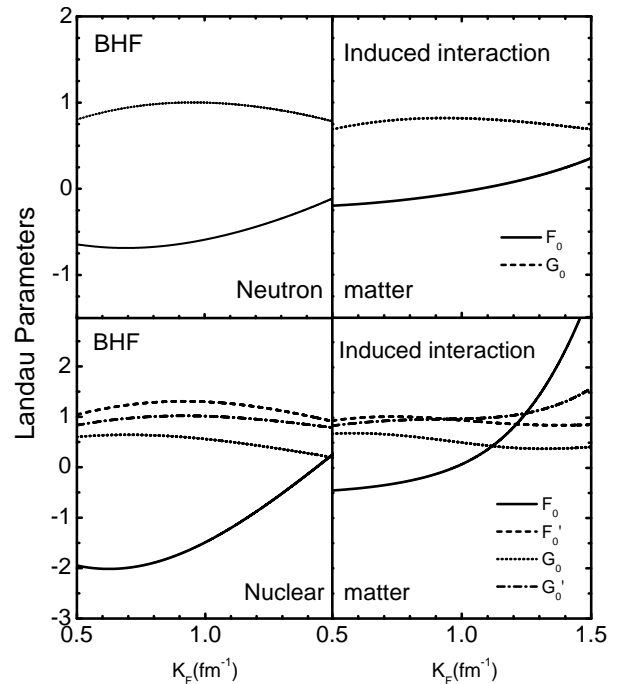

FIG. 3: Landau parameters of pure neutron matter and nuclear matter.

C. Bubble series

The vertex insertions dressing the bubbles must be treated on different footing than the external ones (diagram (c) of Fig. 1). To sum up the RPA bubble series

with the G-matrix requires to solve the Bethe-Salpeter equation, what is a prohibitive task. Using the Landau parameters corresponds to the Landau limit (zero momentum-energy transfer around the Fermi surface), which is a quite reasonable approximation. In this case the RPA summation of the ph interaction turns out to be algebraic and, expressed in term of the dressed bubble, it is written as

$$\Lambda(q)_{ST} = \frac{\Lambda^0(q)}{1 + \Lambda^0(q)\mathcal{L}_{ST}}, \tag{11}$$

where $\mathcal{L}_{ST}$ are the Landau parameters, whose components are commonly denoted by: $\mathcal{L}_{00} = F$, $\mathcal{L}_{01} = F'$, $\mathcal{L}_{10} = G$, $\mathcal{L}_{11} = G'$. In this expression we clearly see how the induced interaction prevents any divergence to occur since $|\Lambda\mathcal{L}| \leq |\mathcal{L}| \leq 1$. Replacing in Eq. (5) the bare bubble Λ^0 with the dressed bubble Λ we get the full screening interaction used in the calculation.

IV. RESULTS

The G-matrix is generated from a selfconsistent BHF calculation with the continuous choice [7]. The Argonne V18 two body force [11] is adopted as the input bare interaction plus a microscopic model for the three body force based on meson exchange with intermediate excitation of nucleon resonances (Delta, Roper, and nucleon-antinucleon) [23]. The calculation also provides the self-energy from which we extract the sp spectrum and the Z-factors. Based on the same framework is also the Skyrme-like fit of the BHF energy functional used to calculate the Landau parameters [20].

A. Screening interaction

In this paper we only focus on the $^1S_0$ pairing interaction in the two extreme situations of pure neutron mat-

ter and symmetric nuclear matter. We keep for a further investigation the consideration of asymmetric nuclear matter with the purpose of studying the transition from the screening regime in pure neutron matter to the antiscreening regime in symmetric nuclear matter.

Let us start with neutron matter. In this case the screening interaction can be decomposed in two terms: S=0 density fluctuation and S=1 spin density fluctuation. The two modes have opposite effect: the former one is attractive, the latter is repulsive. From Eq. (5), we can write

$$\mathcal{V}_1 = \frac{1}{2}\Lambda(q)_0 G_0^{ph}G_0^{ph} - \frac{3}{2}\Lambda(q)_1 G_1^{ph}G_1^{ph}. \tag{12}$$

The factor 3 is due to the multiplicity of the spin mode. For the following discussion we should notice that Λ is negative. In turn, each G-matrix can be expressed as a superposition of G-matrices, projected onto two particle states (see Eq. (6)),

$$G_0^{ph} = \frac{1}{2}(-G_0 - 3G_1), \tag{13}$$

$$G_1^{ph} = \frac{1}{2}(G_0 - G_1). \tag{14}$$

Since G_0 is attractive and G_1 is repulsive, assuming their magnitude to be comparable, we get $G_0 \approx G_1$ [24] and the multiplicity plays the main role in establishing the dominance of the spin density mode over the density mode. In the latter calculation G_0 and G_1 are only roughly comparable, as it can bee seen in Fig.4, nevertheless the conclusion is still valid. A quenching of $^1S_0$ pairing in neutron matter is to be expected, a result established long ago [25] and confirmed by many calculations in various approximations (see [3] and references therein).

In nuclear matter the situation could be quite different since the isospin fluctuations also come into play. The screening interaction now is split according to Eq. (5) as follows

$$\mathcal{V}_1 = \frac{1}{4}(\Lambda(q)_{00}G_{00}^{ph}G_{00}^{ph} + \Lambda(q)_{01}G_{01}^{ph}G_{01}^{ph}) - \frac{3}{4}(\Lambda(q)_{10}G_{10}^{ph}G_{10}^{ph} + \Lambda(q)_{11}G_{11}^{ph}G_{11}^{ph}).$$

The various contributions are plotted in Fig. 4 in terms of pp states, the individual ph contributions are expressed as

$$G_{00}^{ph} = \frac{1}{4}(G_{00} + 3G_{10} + 3G_{01} + 9G_{11}), \tag{15}$$

$$G_{10}^{ph} = \frac{1}{4}(-G_{00} - 3G_{10} + G_{01} + 3G_{11}), \tag{16}$$

$$G_{01}^{ph} = \frac{1}{4}(-G_{00} - 3G_{10} + G_{01} + 3G_{11}), \tag{17}$$

$$G_{11}^{ph} = \frac{1}{4}(G_{00} - G_{10} - G_{01} + G_{11}). \tag{18}$$

In nuclear matter the pp G-matrix elements are dominated by the deuteron channel ($^3SD_1$ coupled pp channel), which is very attractive and therefore it reinforces the density mode and weakens the spin mode. In other words, the main isospin effect is to reverse the role of the medium, i.e. antiscreening instead of screening. In previous papers this effect has been discussed in terms of proton-proton ph screening against neutron-neutron ph screening in the neutron-neutron $^1S_0$ channel [26]. The latter gives repulsion the former attraction. At variance with Ref. [24] the proton-proton ph screening is stronger than neutron-neutron ph screening. This effect

is to be traced back to stronger in medium renormalization of the force in the $T = 0$ channel than in the $T = 0$ one. Antiscreening is the overall effect.

In Fig. 5 we plot the full pairing interaction in the three approximations used in the calculation of the energy gap. In nuclear matter, as we discussed before, the screening effects in fact reinforce the attractive strength of the bare interaction. The main effect appears already at the one bubble level. The deviation from the bare interaction increases at lower density. At $k_F = 0.6$ fm^{-1} the enhancement is from -13 $MeV \cdot fm^3$ to -27 $MeV \cdot fm^3$. This is a huge variation which could entail a large increase of the gap because it is exponentially depending on the interaction. But at such a density the pair correlations are rather weak and thus we do not expect any large increase of the gap. In the density domain of the maximum gap the enhancement is much smaller and again we do not expect any dramatic change in the gap magnitude as an effect of the antiscreening. In Ref. [4], an improper coupling of the ph states in the mixed representation prevented the cancelation among different ph excitations to occur with the effect of producing a more pronounced antiscreening.

In neutron matter the situation is the other way round. The screening is repulsive, and small in the full RPA calculation, but still enough to produce a sizeable quenching of the pairing gap. These predictions confirm at least at qualitative level the corresponding results obtained with the Gogny force [4].
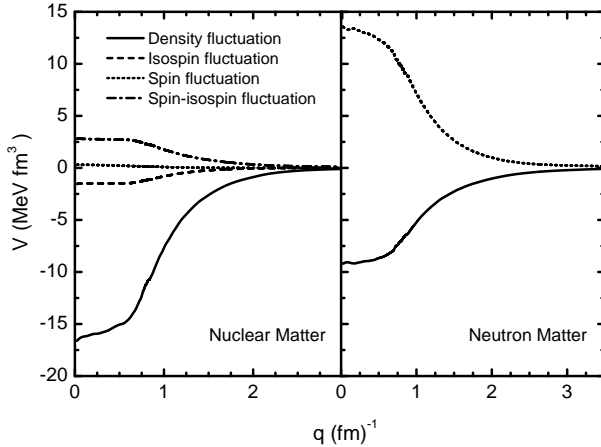

FIG. 4: Individual components of the ph residual interaction.

B. Pairing gap

The present calculation is focussed on the $^1S_0$ neutron-neutron (or proton-proton) pairing. One can distinguish the bare interaction which is responsible for the pairing between the two particles in the $^1S_0$ state, from the screening interaction induced by the surrounding particles. Therefore the interaction, projected onto
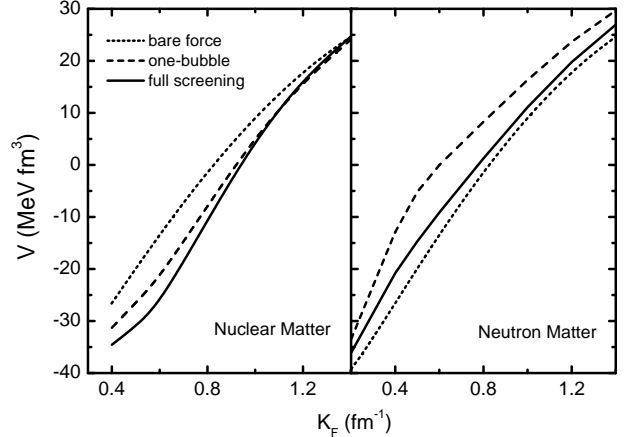

FIG. 5: Pairing interaction.

the $^1S_0$ channel, can be cast as follows

$$< k|\mathcal{V}|k' > = \int \frac{d\Omega}{4\pi} [\mathcal{V}_0(\vec{k}, \vec{k}') + \mathcal{V}_1(|\vec{k} - \vec{k}'|)]. \quad (19)$$

As bare force we use Argonne $V18$, the same as for calculating the G-matrix and the selfenergy. The BCS energy gap in the $^1S_0$ channel is practically independent of the adopted bare force, since in fact all realistic interactions reproduce the phase shifts of free NN scattering.

We solved the gap equation in the form of Eq. (4). In order to disentangle the screening effects from the selfenergy corrections, we first assume Z=1 and free sp spectrum. The results are plotted in Fig. 6 (upper left panel). In neutron matter the screening effect is small and just reduces the gap by 10% in the peak region. At variance with previous calculations existing in the literature [3] the full RPA screening is much less effective than the one bubble approximation because of the stronger renormalization of the spin fluctuations vs the density fluctuations in the induced interaction. However this finding confirms the preceding predictions with Gogny force (see Fig.8 of Ref. [4]).

In nuclear matter, due to the antiscreening effect we discussed earlier, the magnitude of the gap variation is the other way around and much more sizeable: the gap rises up from 3 MeV to 5 MeV for Fermi momentum $k_F = 0.8$ fm^{-1}. This is displayed in Fig. 6 (lower left panel).

There are two kinds of selfenergy effects: dispersive effect and Fermi surface depletion. Both are calculated taking into account the selfenergy corrections at the second order of G-matrix (rearrangement terms). The first one is a correction to the sp spectrum in the energy denominator. Usually it entails a reduction of the pairing gap since the effective mass, beyond BHF approximation, is less than the unity (the effective mass is the combination of the e-mass and the k-mass [7]). But at very low density the effective mass is larger than unity [27] and it reduces the quenching rate of the gap due to the interaction. This effect can be seen in the low density side of the neutron gap with Σ_{total} (upper right panel). Additional strong reduction is due to the de-

pletion of the Fermi surface which hinders transitions around the Fermi surface. The maximum gap in a complete calculation is $1.5 - 2 \ MeV$ at $k_F \approx 0.8 \ fm^{-1}$.

In nuclear matter the selfenergy effects are much stronger already at moderately low density, as it has to be expected, and the peak value shifts down to very low density $k_F \approx 0.5 - 0.6 \ fm^{-1}$. The Z-factor plays the major role: it quenches from 0.84 in neutron matter to 0.68 in nuclear matter at $k_F = 0.8 \ fm^{-1}$. But the magnitude is about 0.5 MeV less than the value with only bare interaction. Therefore we can conclude that a strong cancelation occurs as soon as vertex corrections and self energy effects are simultaneously included in the gap equation. But this happens only in nuclear matter as an effect of antiscreening. We will come back to this point below.
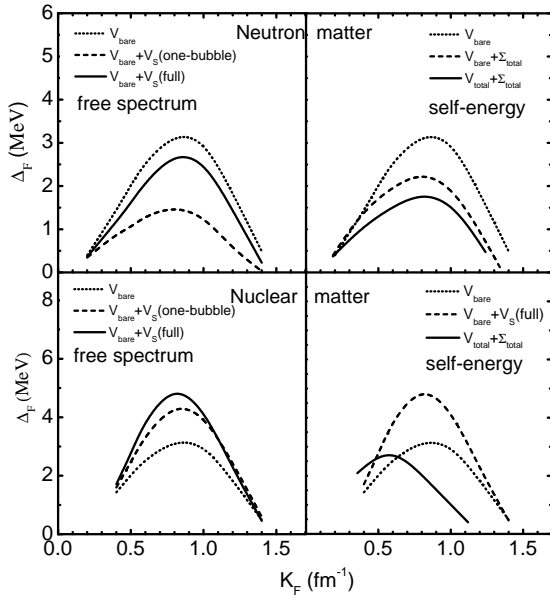

FIG. 6: Pairing gap in the $^1S_0$ channel for pure neutron matter (upper figure) and nuclear matter (lower figure).

V. DISCUSSION AND CONCLUSIONS

In this paper an exhaustive treatment of the $^1S_0$ pairing in nuclear and neutron matter has been reported. The medium polarization effects on the interaction and the selfenergy corrections to the mean field, both developed in the framework of the Brueckner theory, have been included in the solution of the gap equation.

Within the pure mean field approximation [28] the $^1S_0$ gap is not affected by the medium, either nuclear or neutron matter. So far the medium effects have not been considered in the case of nuclear matter except in Ref. [4]. The vertex corrections due to neutron matter all give a reduction of the pairing, the magnitude depending on the adopted approximation [3]. The expla-

nation relies on the competition of the attractive density excitations against the repulsive spin density excitations. The present calculation, based on G-matrix, also predicts a large quenching in agreement with almost all previous predictions, but only at the one bubble level. In the most complete calculation (full RPA) the quenching is largely reduced in apparent agreement with a recent Monte Carlo calculation [18]. But the inclusion of selfenergy effects definitely results in a large suppression as expected from basic properties of a strongly correlated many body system (see Introduction).

In the case of nuclear matter the most remarkable result is the antiscreening effect of the medium polarization. In fact in nuclear matter isospin modes arise that reverts the competition between the attractive density modes and the repulsive spin-density modes due to the presence of isospin modes. The argument addressed in Ref. [24] that the p-n (T=0) interaction is small compared to the n-n (T=1) is based on the vacuum scattering T-matrix and does not consider the strong medium renormalization of G-matrix, which inverts the strength of the two channels. However the enhancement of the gap to almost 5 MeV is almost completely suppressed by the strong correlation effects on the selfenergy. But, even a small variation of the force strength implies a large variation of the gap. These effects also push to lower density the peak value of the gap.

Calculations of the pairing gap in a nuclear environment have been reported in a series of papers for the case of nuclei [5]. Their main finding is that the induced interaction arising from the surface vibrations is responsible for large part of the experimental gap. This result can be considered as the counterpart for finite nuclei of the antiscreening effect due to the medium polarization in nuclear matter. But the selfenergy effects completely compensate the gap enhancement and, in the end, the full medium effect do not change significantly the gap with bare interaction. This result turns out to be not a big surprise, since some calculations show that the gap with Gogny interaction is consistent with the observed gaps in nuclei [29].

At this point two aspect are worth to be developed further. The first one is the study of pairing in the transition from symmetric nuclear matter to neutron matter; the second one is the investigation of pp channels so far neglected, since there the bare interaction is repulsive. The pairing could exist as an induced effect of the environment.

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